Split Two-Higgs-Doublet Model and Neutrino Condensation

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We split the two-Higgs-doublet model by assuming very different vevs for the two doublets: the vev is at weak scale (174 GeV) for the doublet Φ_1 and at neutrino-mass scale ($10^{-2} \sim 10^{-3}$ eV) for the doublet Φ_2. Φ_1 is responsible for giving masses to all fermions except neutrinos; while Φ_2 is responsible for giving neutrino masses through its tiny vev without introducing see-saw mechanism. Among the predicted five physical scalars H, h, A^0 and $H^\pm$, the CP-even scalar h is as light as $10^{-2} \sim 10^{-3}$ eV while others are at weak scale. We identify h as the cosmic dark energy field and the other CP-even scalar H as the Standard Model Higgs boson; while the CP-odd A^0 and the charged $H^\pm$ are the exotic scalars to be discovered at future colliders. Also we demonstrate a possible dynamical origin for the doublet Φ_2 from neutrino condensation caused by some unknown dynamics.

12.60.Fr,12.60.Rc,14.60.Pq

Introduction: The two-Higgs-doublet model (2HDM) [1] is the most naive extension of the Standard Model (SM). However, such an extension seems to bring more heat than light: it plainly introduces several more scalar particles without solving any problems in the SM. In this work we reform this model in order to explain the tiny neutrino masses without see-saw mechanism and provide a candidate field for cosmic dark energy. For this purpose we split this model by assuming very different vevs for the two doublets: the vev is at weak scale (174 GeV) for the doublet Φ_1 and at neutrino-mass scale ($10^{-2} \sim 10^{-3}$ eV) for the doublet Φ_2. We assume that Φ_1 is responsible for giving masses to all fermions except neutrinos; while Φ_2 is responsible for giving neutrino masses through its tiny vev without introducing see-saw mechanism. Among the predicted five physical scalars H, h, A^0 and $H^\pm$, the CP-even scalar h is as light as $10^{-2} \sim 10^{-3}$ eV (in the following we assume 10^{-3} eV for example) while others are at weak scale. We identify h as the cosmic dark energy field and the other CP-even H as the SM Higgs boson; while the CP-odd A^0 and the charged $H^\pm$ are the exotic scalars to be discovered at future colliders.

Our motivations for splitting the 2HDM are then quite clear:

- By identifying h as the cosmic dark energy field, we try to give an explanation for dark energy (although we know we may be quite far from the true story). As is well known, the dark energy seems to be a great mystery in today's physics and cosmology. The precise cosmological measurements, such as the Wilkinson Microwave Anisotropy Probe (WMAP) measurements [2], indicate that the major content of today's universe is the weird dark energy. So far we lack the understanding about the nature of dark energy, although some phenomenological approaches have been proposed [3].

- By giving neutrino masses through the tiny vev of Φ_2, we try to *'understand'* the tiny neutrino masses without introducing see-saw mechanism (although we know the see-saw mechanism is quite elegant). It is also well known that among the elementary particles the neutrinos are quite special species due to their extremely small masses. Without see-saw mechanism, the Yukawa couplings of neutrinos with the SM Higgs doublet are extremely small, which is hard to understand. We may speculate that the origin of neutrino masses are different from other fermions, i.e., the neutrino masses are not from the Yukawa couplings of the SM Higgs doublet (there might be some symmetry to forbid the neutrino couplings with the SM Higgs doublet), and, instead, they are from the Yukawa couplings of a new scalar doublet which has a tiny vev.

- We try to relate neutrino mass generation with dark energy puzzle. The neutrino mass scale is seemingly near or coincident with the cosmic dark energy scale (the dark energy density is $\sim (10^{-3} eV)^4$ and thus the dark energy scale is $\sim 10^{-3} eV$). Such an seeming coincidence has already stimulated some speculations on the possible relation between neutrino and dark energy [4].

While the dynamical origin of the scalar doublet Φ_1 with a vev at weak scale may be something like technicolor, we propose that the scalar doublet Φ_2 with a tiny vev at the neutrino-mass scale may be from neutrino condensation. Similar to the idea of top-quark condensation [5], we assume that a four-fermion interaction for the third-family neutrino is induced at some high energy scale (say TeV) from some unknown new dynamics (like top color [6]) which is strong enough to cause neutrino

condensation [1].

Split two-Higgs-doublet model: We introduce two scalar doublets Φ_1 and Φ_2 as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \text{Re}\phi_1^0 + v_1 + i\text{Im}\phi_1^0 \end{pmatrix}, \tag{1}$$

$$\Phi_2 = \begin{pmatrix} \phi_2^+ \\ \text{Re}\phi_2^0 + v_2 + i\text{Im}\phi_2^0 \end{pmatrix}. \tag{2}$$

We split the two vevs as

$$v_1 \sim 174 \text{ GeV}, \quad v_2 \sim 10^{-3} \text{ eV}. \tag{3}$$

By assumption, Φ_1 is responsible for giving masses to all fermions except neutrinos; while Φ_2 is responsible for giving neutrino masses through its tiny vev. Of course, both of them contribute to W boson mass:

$$m_W^2 = g^2(v_1^2 + v_2^2)/2 \approx g^2 v_1^2/2. \tag{4}$$

We assume CP conservation and discrete symmetry $\Phi_1 \to -\Phi_1$ for the potential of the scalars. Then the general potential takes the form [1]

$$\begin{aligned} V(\Phi_1, \Phi_2) = {} & \lambda_1(\Phi_1^\dagger\Phi_1 - v_1^2)^2 + \lambda_2(\Phi_2^\dagger\Phi_2 - v_2^2)^2 \\ & +\lambda_3\left[(\Phi_1^\dagger\Phi_1 - v_1^2) + (\Phi_2^\dagger\Phi_2 - v_2^2)\right]^2 \\ & +\lambda_4\left[(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) - (\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1)\right] \\ & +\lambda_6\left[\text{Im}(\Phi_1^\dagger\Phi_2)\right]^2, \end{aligned} \tag{5}$$

where λ_i are non-negative real parameters. Note that since we assumed the exact discrete symmetry $\Phi_1 \to -\Phi_1$, we dropped the term

$$\lambda_5\left[\text{Re } (\Phi_1^\dagger\Phi_2) - v_1 v_2\right]^2, \tag{6}$$

which softly breaks such a discrete symmetry.

It should be pointed out that since the two vevs of the potential in eq.(5) are splitted by many orders, in order to make it stable under radiative corrections we either need an extreme fine tuning as that occuring in GUTs or introduce some fancy symmerty to stablize it.

After the diagonalization of the mass-square matrices in Eq.(5), we obtain eight mass eigenstates: H, h, $H^\pm$, $G^\pm$, A^0 and G^0, among which H is identified as the SM physical Higgs boson, and $G^\pm$ and G^0 are massless Goldstone boson eaten by W and Z gauge bosons.

The two charged scalars $H^\pm$ and $G^\pm$ are obtained by

$$G^\pm = \Phi_1^\pm \cos\beta + \Phi_2^\pm \sin\beta \approx \Phi_1^\pm, \tag{7}$$

$$H^\pm = -\Phi_1^\pm \sin\beta + \Phi_2^\pm \cos\beta \approx \Phi_2^\pm, \tag{8}$$

while the two CP-odd neutral scalars A^0 and G^0 are obtained by

$$G^0 = \sqrt{2}(\text{Im}\Phi_1^0 \cos\beta + \text{Im}\Phi_2^0 \sin\beta) \approx \sqrt{2}\text{Im}\Phi_1^0, \tag{9}$$

$$A^0 = \sqrt{2}(-\text{Im}\Phi_1^0 \sin\beta + \text{Im}\Phi_2^0 \cos\beta) \approx \sqrt{2}\text{Im}\Phi_2^0, \tag{10}$$

where the mixing angle β is very small since $\tan\beta = v_2/v_1$.

The masses of $H^\pm$ and A^0 are given by

$$m_{H^\pm}^2 = \lambda_4(v_1^2 + v_2^2) \approx \lambda_4 v_1^2, \tag{11}$$

$$m_{A^0}^2 = \lambda_6(v_1^2 + v_2^2) \approx \lambda_6 v_1^2. \tag{12}$$

So, both $m_{H^\pm}$ and m_{A^0} should be at weak scale given that $\lambda_i \sim \mathcal{O}(1)$.

The two CP-even scalars H and h are obtained by

$$\begin{aligned} H &= \sqrt{2}\left[(\text{Re } \Phi_1^0 - v_1)\cos\alpha + (\text{Re } \Phi_2^0 - v_2)\sin\alpha\right] \\ &\approx \sqrt{2}(\text{Re } \Phi_1^0 - v_1), \end{aligned} \tag{13}$$

$$\begin{aligned} h &= \sqrt{2}\left[-(\text{Re } \Phi_1^0 - v_1)\sin\alpha + (\text{Re } \Phi_2^0 - v_2)\cos\alpha\right] \\ &\approx \sqrt{2}(\text{Re } \Phi_2^0 - v_2). \end{aligned} \tag{14}$$

Here the mixing angle α is also very small since

$$\tan(2\alpha) = \frac{2\mathcal{M}_{12}}{\mathcal{M}_{11} - \mathcal{M}_{22}} \sim v_2/v_1, \tag{15}$$

where $\mathcal{M}_{ij}$ are elements of the mass-square matrix

$$\mathcal{M} = \begin{pmatrix} 4(\lambda_1 + \lambda_3)v_1^2 & 4\lambda_3 v_1 v_2 \\ 4\lambda_3 v_1 v_2 & 4(\lambda_2 + \lambda_3)v_2^2 \end{pmatrix}. \tag{16}$$

The masses of H and h are given by

$$m_H^2 \approx \mathcal{M}_{11} = 4(\lambda_1 + \lambda_3)v_1^2, \tag{17}$$

$$m_h^2 \approx \mathcal{M}_{22} - \frac{\mathcal{M}_{12}^2}{\mathcal{M}_{11}} = 4\left[(\lambda_2 + \lambda_3) - \frac{\lambda_3^2}{(\lambda_1 + \lambda_3)}\right]v_2^2. \tag{18}$$

Therefore, among the five physical scalars, the masses of $H^\pm$, A^0 and H are at the weak scale since they are proportional to $v_1 \approx 174$ GeV, while the mass of h is at the neutrino mass scale since it is proportional to $v_2 \sim 10^{-3}$ eV. The two mixing angles α and β are very small since they are proportional to v_2/v_1.

Due to the negligibly small mixing angles α and β, the properties of these scalars are approximately like this:

(i) For H: All its couplings are the same as in the SM.

(ii) For $H^\pm$: They have Yukawa couplings only to $\bar{\ell}_L \nu_R^\ell$ ($\ell = e, \mu, \tau$), but the coupling strength is at a natural order, say $\mathcal{O}(1)$. Their gauge couplings like $H^+H^-\gamma$ and H^+H^-Z are the same as in the usual 2HDM.

[1] Note that in the literature the neutrino condensation was once proposed as a dynamical electroweak symmetry breaking mechanism [7], which can generate tiny neutrino mass by incorporating the see-saw mechanism.

(iii) For A^0: It has Yukawa couplings only to neutrino pairs. It has gauge couplings like ZhA^0 as in the usual 2HDM.

(iv) For h: It has Yukawa couplings only to neutrino pairs. Its gauge couplings to W^+W^- and ZZ are very weak (proportional to v_2).

Therefore, the experimental constraints on these scalars are:

(1) For H: all direct and indirect experimental constraints are the same as for the SM Higgs boson.

(2) For $H^\pm$ and A^0: the experimental constraints are similar as in the usual 2HDM except for the invalid constraints from various B-decays (such as $b \to s\gamma$). For example, from the unobservation of $e^+e^- \to H^+H^-$ at LEP II, $H^\pm$ should be heavier than about 100 GeV and thus λ_4 in Eq. (11) cannot be too small.

(3) For the ultra-light h: stringent constraints both from particle physics and from astrophysics are derived from its interactions with either photons, electrons or nucleons [8]. For example, stringent constraints may come from positronium decays, meson decays, quarkonium decays or nuclear transitions. In our model, fortunately, these constraints can be avoided or become quite weak since the coupling of h with photons, electrons or nucleons are suppressed by $v_2/v_1 \sim 10^{-14}$ at tree-level. The most dangerous constraints may come from the invisible Z decays. For example, from the three-body decay $Z \to h(A^0)^* \to h\nu\bar{\nu}$, some lower mass bound (say TeV) may be set on A^0.

Note that the light scalar h has a Yukawa coupling of order one to neutrinos (left and right handed). Thus, in addition to left-handed neutrinos, the right-handed neutrinos and the scalar h also appear in the thermal equilibrium in the early universe. This means that the effective total number of neutrino species (N_ν) is about 6.5 instead of the standard number of 3. We should check if this is allowed by cosmology and astrophysics (right-handed neutrinos are subject to no constraints from particle physics experiments such as LEP experiments since they are gauge singlet and have no gauge couplings). Firstly, we note that stringent constraint on N_ν has been derived from the standard BBN [9]. However, as discussed in [10], such BBN contstraint can be relaxed since it is obtained under the assumption that the chemical potential for the background neutrinos is negligible. The relaxed 2σ bound from the combined analysis of BBN, CMB and supernova data is $N_\nu < 7$ [10]. Secondly, we should seriously and specially consider the constraints from SN1987A since the additional neutrino species can speed up the cooling-down of its core. It is well known that there exist some uncertainty for the estimation of the total binding energy as well as the effective temperature of the supernova. As analysed in [11], the constraint is $N_\nu < 6.7$ (the Eq.8 in [11]).

It should also be noted that since the light scalar h can couple to ordinary matter through the mixing with the SM Higgs boson H (suppressed by $v_2/v_1 \sim 10^{-14}$) or radiatively through loops involving left- and right-handed neutrinos plus weak gauge bosons (also suppressed by v_2/v_1 because right-handed neutrinos do not have any gauge interactions and a neutrino mass insertion is needed in order to couple to gauge bosons), it is necessary to check whether the gravitational equivalence principle is still respected in our case. As analysed in [12], it will be consistent with tests of the gravitational inverse square law as well as the equivalence principle as long as the scalar is not too light or its Yukawa couplings to ordinary matter are sufficiently weak. Following the analysis in [12], we checked that our case is marginally in the region allowed by the gravitational inverse square law and the equivalence principle.

Neutrino Condensation: Now we try to provide an explanation for the generation of Φ_2 by neutrino condensation. We assume that the third-family leptons have an effective four-fermion interaction, which may be generated at some high energy scale Λ above TeV from some unknown new dynamics. The underlying new dynamics, which is not specified here, might be some non-abelian gauge interaction spontaneously broken at some higher scale. So such new dynamics causes negligible effects to the electroweak physics of the third-family leptons.

The four-fermion effective interaction for the third-family leptons takes the form at the energy scale Λ

$$G\left(\bar{\Psi}_L \nu_R\right)\left(\bar{\nu}_R \Psi_L\right), \tag{19}$$

where G is the coupling constant, running with the energy scale. Ψ_L is the doublet of the left-handed third-family lepton fields and ν_R is the right-handed tau-neutrino. When $G\Lambda^2 \gg 1$, the tau-neutrinos condensate and the condensation effects can be incorporated by introducing an auxiliary scalar field Φ_2 into the Lagrangian

$$-M_0\sqrt{G}\left[\bar{\Psi}_L \tilde{\Phi}_2 \nu_R + h.c.\right] - M_0^2 \tilde{\Phi}_2^\dagger \tilde{\Phi}_2 \tag{20}$$

where M_0 is an unspecified bare mass parameter and $\tilde{\Phi}_2 = i\tau_2 \Phi_2^*$. Φ_2 and $\tilde{\Phi}_2$ take the form

$$\Phi_2 = \frac{\sqrt{G}}{2M_0}\begin{pmatrix} \bar{\tau}(1+\gamma_5)\nu \\ -\bar{\nu}(1+\gamma_5)\nu \end{pmatrix} = \frac{\sqrt{G}}{M_0}(i\tau_2)(\bar{\nu}_R\Psi_L)^{\dagger T} \tag{21}$$

$$\tilde{\Phi}_2 = -\frac{\sqrt{G}}{M_0}(\bar{\nu}_R \Psi_L). \tag{22}$$

As the energy scale runs down, such an auxiliary field Φ_2 gets gauge invariant kinematic terms as well as quartic interactions through quantum effects:

$$\begin{aligned} &-M_0\sqrt{G}\left(\bar{\Psi}_L\tilde{\Phi}_2\nu_R + h.c.\right) + Z_{\Phi_2}|D_\mu\Phi_2|^2 \\ &-M_{\Phi_2}^2\Phi_2^\dagger\Phi_2 - \lambda_0\left(\Phi_2^\dagger\Phi_2\right)^2, \end{aligned} \tag{23}$$

where

$$Z_{\Phi_2} = \frac{N_c M_0^2 G}{(4\pi)^2}\ln\left(\frac{\Lambda^2}{\mu^2}\right) \tag{24}$$

$$M_{\Phi_2}^2 = M_0^2 - \frac{2N_c M_0^2 G}{(4\pi)^2}(\Lambda^2 - \mu^2) \tag{25}$$

$$\lambda_0 = \frac{N_c(M_0^2 G)^2}{(4\pi)^2}\ln\left(\frac{\Lambda^2}{\mu^2}\right). \tag{26}$$

Here N_C is the 'color' number of tau-neutrino in the unspecified new dynamics. Redefining $\sqrt{Z_{\phi_2}}\Phi_2$ as Φ_2 and $g_t = M_0\sqrt{G}/\sqrt{Z_{\Phi_2}}$, $m^2 = M_{\Phi_2}^2/Z_{\Phi_2}$, $\lambda = \lambda_0/Z_{\Phi_2}^2$, we obtain the Lagrangian

$$\begin{aligned} &-g_t\left(\bar{\Psi}_L\tilde{\Phi}_2\nu_R + h.c.\right) + |D_\mu\Phi_2|^2 \\ &-m^2\left(\Phi_2^\dagger\Phi_2\right) - \lambda\left(\Phi_2^\dagger\Phi_2\right)^2 \end{aligned} \tag{27}$$

For $\mu \ll \Lambda$, we have $m^2 < 0$ when $G \geq G_{crit} = 8\pi^2/(N_c\Lambda^2)$, which leads to spontaneous breaking of electroweak symmetry with Φ_2 developing a vev $v_2 = \sqrt{-m^2/2\lambda}$.

The masses of tau-neutrino is given by

$$m_\nu^2 = g_t^2 v_2^2 = \frac{16\pi^2 v_2^2}{N_c\ln(\Lambda^2/m_\nu^2)}, \tag{28}$$

which is consistent with $m_\nu \sim v_2 \sim 10^{-3}$ eV for $\Lambda \sim$ TeV.

Collider and Cosmological Consequences: First, we briefly discuss the phenomenology at the colliders, LHC (CERN Large Hadron Collider) and ILC (International Linear Collider). Since the phenomenology of H is approximately the same as the SM Higgs boson, we focus on the phenomenology of $H^\pm$, A^0 and h.

(1) Since h is as light as neutrino and couples with neutrinos, its dominant decay mode is $h \to \nu\bar{\nu}$. Thus it just escapes detection at colliders although it can be produced through $e^+e^- \to Z^* \to hA^0$ at the ILC.

(2) At the ILC A^0 can be produced through $e^+e^- \to Z^* \to hA^0$ followed by the decay $A^0 \to Zh$, leading to the signature of two lepton or two jets plus missing energy.

(3) At the ILC $H^\pm$ can be pair produced not only through the usual s-channel process $e^+e^- \to Z^*$, $\gamma^* \to H^+H^-$ but also through the t-channel process via exchanging a ν_R. Due to such an additional t-channel process, the production rate will be different from the usual 2HDM prediction. The decay modes $H^- \to \ell_L\bar{\nu}_R$ ($\ell = e,\mu$) will give good signatures. Since $H^\pm$ almost do not couple to quarks, they cannot be produced at LHC collider through the subprocess like $gb \to tH^-$.

Now we turn to the cosmological consequences. Since the CP-even physical scalar h is an ultra-light scalar particle, at the order of 10^{-3} eV, we may interpret this scalar field as the dark energy field. Since the corresponding vev v_2 is at the order of 10^{-3} eV, the magnitude of the vacuum energy (dark energy) should be naturally around the order of $(10^{-3}\ eV)^4$. This means that in fact we assume something that sends to zero the potential energy related to the usual Weinberg-Salam Higgs doublet while leaving the potential energy related to the light scalar h unbalanced.

Since h may decay into $\nu\bar{\nu}$, we wonder if it is still present in today's universe as a viable dark energy field. As is well known, there is now a neutrino microwave background at 1.9 K which corresponds to $\sim 10^{-4}$ eV. The scalar h and the neutrinos can reach thermal equilibrium if the mass difference between h its decay products $\nu\bar{\nu}$ is at the same order as the neutrino background temperature. Assuming the contribution of chemical potential is at the order of unity, the ratio for n_h and n_ν in equilibrium is proportional to $exp(-\Delta Q/T)$, where ΔQ is the released energy in the decay process.

Note that just like some authors in [4], we may assume the dark fluid to be the sum of the scalar potential of h and the energy density in neutrino masses. In this case, since the scalar potential evolves with energy scale (time), the minimum of the resulting potential would evolve in time and also one could have mass varying neutrinos.

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